

# OBLIQUE RADIATIVE CORRECTIONS IN THE VECTOR CONDENSATE MODEL OF ELECTROWEAK INTERACTIONS

G. Cynolter<sup>1</sup>, E. Lendvai<sup>2</sup> and G. Pócsik<sup>3</sup>

Institute for Theoretical Physics, Eotvos Lorand University,  
Budapest

## Abstract

Oblique radiative corrections are calculated to the parameter  $S$  in a version of the standard model where the Higgs doublet is replaced by a doublet of vector bosons and the gauge symmetry is broken dynamically. We show that to each momentum scale there exists a domain of the masses of charged and neutral vector bosons where  $S$  is compatible with the experiments. At a scale of 1 TeV this requires vector boson masses of at least  $m_0 \approx 400\text{--}550$  GeV,  $m_+ \approx 200\text{--}350$  GeV.

The status of the standard model is very good, but it is not clear whether the Higgs mechanism is responsible for the origin of the mass. Therefore, alternative realizations of the electroweak symmetry breaking are also important [1].

Recently, we have proposed a model [2,3] of electroweak symmetry breaking where the standard model scalar doublet is replaced by a doublet of vector fields,

$$B_\mu = \begin{pmatrix} B_\mu^{(+)} \\ B_\mu^{(0)} \end{pmatrix}, \quad (1)$$

and  $B_\mu^{(0)}$  forms a condensate

$$\langle B_\mu^{(0)+} B_\nu^{(0)} \rangle_0 = g_{\mu\nu} d, \quad d \neq 0. \quad (2)$$

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<sup>1</sup>E-mail: cyn@hal9000.elte.hu

<sup>2</sup>E-mail: lendvai@hal9000.elte.hu

<sup>3</sup>E-mail: pocsik@ludens.elte.hu

$B_\mu$  is coupled to the  $SU(2) \times U(1)$  gauge fields and to itself by a gauge invariant Lagrangian

$$L = -\frac{1}{2} (D_\mu B_\nu - D_\nu B_\mu)^+ (D^\mu B^\nu - D^\nu B^\mu) - \lambda (B_\nu^+ B^\nu)^2, \lambda > 0, \quad (3)$$

with  $D_\mu$  the usual covariant derivative. Fermion and gauge field couplings are standard.

Breaking the symmetry in (3) by (2) gives rise to the usual mass terms for the gauge fields and makes also the B particles massive.  $\sqrt{-6d} \simeq 246\text{GeV}$  plays the role of the vacuum expectation value of the scalar field [2] and  $\rho_{tree} = 1$ .

Fermion masses are considered as manifestations of the condensation (2) and derived from the Lagrangian

$$g_{ij}^u (\bar{\Psi}_{Li} B_\nu^C u_{Rj}) B_\nu^{(0)} + g_{ij}^d (\bar{\Psi}_{Li} B_\nu d_{Rj}) B_\nu^{(0)+} + h.c.,$$

$$\Psi_{Li} = \begin{pmatrix} u_i \\ d_i \end{pmatrix}_L B_\nu^C = \begin{pmatrix} B_\nu^{(0)+} \\ -B_\nu^{(+)+} \end{pmatrix} \quad (4)$$

(and similarly for leptons) [3]. (4) allows the Kobayashi–Maskawa mechanism, too. (3) and (4) fix the interactions of B particles. Producing B particles in  $e^+e^-$  annihilation has been investigated in ref. [2].

The model can be considered as a low energy effective model valid up to a scale  $\Lambda$ ,  $\Lambda \approx 2.6\text{TeV}$ . We have shown [3] that oblique radiative corrections due to B-loops give arbitrarily small contributions to the  $\rho$  parameter if  $B^{+,0}$  masses are suitably chosen, and  $\Lambda$  remains unrestricted.

In the present paper one-loop oblique radiative corrections are calculated to the parameter S, one of the three parameters [4] constrained by precision experiments. It is shown that the B particles must be heavy and to each  $\Lambda$  there exists a domain of  $B^{+,0}$  masses where the model remains valid. For B masses which are very large compared to a fixed  $\Lambda$ , B-loops contribute negligibly to S. At a scale of 1 TeV the threshold is about  $m_+ \approx 200\text{--}350$  GeV,  $m_0 \approx 400\text{--}550$  GeV.

The parameter S [4] defined by

$$\alpha S = 4e^2 (\bar{\Pi}'_{ZZ}(0) - (c^2 - s^2) \bar{\Pi}'_{ZA}(0) - s^2 c^2 \bar{\Pi}'_{AA}(0)),$$

$$\bar{\Pi}'_{ik}(0) = \frac{d}{dq^2} \bar{\Pi}'_{ik}(q^2)|_{q^2=0} \quad (5)$$

with  $c = \cos \theta_w$ ,  $s = \sin \theta_w$  is calculated in one B-loop order.  $\overline{\Pi}'_{ik}(q^2)$  is expressed by the  $g_{\mu\nu}$  terms of the vacuum polarization contributions  $\Pi_{ik}(q^2)$  due to B-loops as

$$\Pi_{ZZ} = \frac{e^2}{s^2 c^2} \overline{\Pi}_{ZZ}, \quad \Pi_{ZA} = \frac{e^2}{s c} \overline{\Pi}_{ZA}, \quad \Pi_{AA} = e^2 \overline{\Pi}_{AA}. \quad (6)$$

To one B-loop order only trilinear interactions of B are essential, from (3) these are as follows

$$\begin{aligned} L(B^0) &= \frac{ig}{2c} \partial^\mu B^{(0)\nu+} (Z_\mu B_\nu^{(0)} - Z_\nu B_\mu^{(0)}) + h.c., \\ L(B^+ B^{++} Z) &= -(c^2 - s^2) \cdot L(B^0 \rightarrow B^+), \\ L(B^+ B^- A) &= ie \partial^\mu B^{(+)\nu+} (A_\mu B_\nu^{(+)} - A_\nu B_\mu^{(+)}) + h.c.. \end{aligned} \quad (7)$$

In a renormalizable model S is finite. In the model based on (3), however, cancellations of  $\Lambda$  powers are not perfect due to nonrenormalizability. We get for the scaled vacuum polarizations  $\overline{\Pi}'_{ik}(0)$ ,

$$\begin{aligned} \overline{\Pi}'_{ZZ}(0) &= \frac{1}{128\pi^2} \left( f\left(\frac{\Lambda^2}{m_0^2}\right) + (c^2 - s^2)^2 f\left(\frac{\Lambda^2}{m_+^2}\right) \right), \\ \overline{\Pi}'_{ZA}(0) &= -\frac{c^2 - s^2}{64\pi^2} f\left(\frac{\Lambda^2}{m_+^2}\right), \\ \overline{\Pi}'_{AA}(0) &= \frac{1}{32\pi^2} f\left(\frac{\Lambda^2}{m_+^2}\right), \end{aligned} \quad (8)$$

where  $m_+(m_0)$  is the physical mass of  $B^+(B^0)$ , and

$$f(x) = 5x - 16 \ln(1+x) + 14 - \frac{17}{1+x} + \frac{3}{(1+x)^2}. \quad (9)$$

For increasing  $m$   $f\left(\frac{\Lambda^2}{m^2}\right)$  decreases through  $f=0$  at  $m = 1.90\Lambda$ , and at  $\frac{\Lambda}{m} \rightarrow 0$   $f \rightarrow 0$  from below.

Substituting (8) into (5) we obtain

$$S = \frac{1}{8\pi} \left[ f\left(\frac{\Lambda^2}{m_0^2}\right) + (3(c^2 - s^2)^2 - 4s^2 c^2) f\left(\frac{\Lambda^2}{m_+^2}\right) \right]. \quad (10)$$

The coefficient of  $f\left(\frac{\Lambda^2}{m_+^2}\right)$  is 0.158 for  $s^2 = 0.231$ .

An analysis of precision experiments shows that  $S_{new} < 0.09(0.23)$  at 90 (95)% C.L. [5] for  $m_H^{ref} = 300\text{GeV}$  and assuming  $m_t = 174\text{GeV}$  (CDF value). Requiring  $S_{new} \geq 0$ , the corresponding constraints are  $S_{new} < 0.38(0.46)$  [5]. Since a Higgs of 300 GeV is absent in the present model, its contribution, 0.063, must be removed. In this way for the contribution of B we have  $S < 0.15(0.29)$  at 90 (95)% C.L. For  $m_+, m_0 \geq 1.90\Lambda$  this is fulfilled, in particular,  $S \rightarrow 0$  for  $\frac{\Lambda}{m_{+,0}} \rightarrow 0$ . One can fulfil the experimental constraints also for a mass(es) lower than  $1.90\Lambda$ . This is shown in Fig.1. for  $\Lambda = 1\text{TeV}$ , where the curves are coming from the 90% and 95% C.L. limits on S for  $S_{new} \geq 0$  (lower two curves) and for  $S \geq 0$  (upper two curves), respectively. Excluded regions are below the curves. It follows that the B particles are heavy; at  $\Lambda = 1\text{TeV}$  the threshold is about  $m_0 \approx 400\text{--}550\text{ GeV}$  and  $m_+ \approx 200\text{--}350\text{ GeV}$ . Since S is invariant multiplying  $\Lambda, m_0, m_+$  by a common factor, Fig.1. determines the allowed regions for scales different from 1 TeV. Higher  $\Lambda$  attracts higher minimum masses. Recently, we have shown [6] that B particles could be seen at high energy linear  $e^+e^-$  colliders up to masses of several hundred GeV's.

In conclusion, in the vector condensate model, to each momentum scale there exists a range of B boson masses where S is compatible with the experimental constraints. Further restrictions are imposed on  $m_{+,0}$  by taking into account the results [3] on the parameter T. While for a fixed  $\Lambda$  and  $m_0$  a large  $m_+$  range exists from S (see Fig.1.), this is tightened by T. For example, at  $\Lambda = 1\text{TeV}$ ,  $m_0 = 400(600)\text{ GeV}$ , the  $m_+$  range allowed by S, T is  $m_+ = 630\text{--}636(846\text{--}868)\text{ GeV}$ . For higher  $\Lambda$  the allowed  $m_+$  region shrinks at the same  $m_0$ .

In general,  $\Lambda$  remains unrestricted and suitable, heavy  $B^{+,0}$  provide small radiative corrections.

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# FIGURE CAPTION

Fig. 1.  $m_+$  vs.  $m_0$  at  $\Lambda = 1\text{TeV}$ . Lower (upper) two curves embody the 90% and 95% C.L. on S for  $S_{new} \geq 0$  ( $S_{new} \gtrless 0$ ). Allowed regions are above the curves.

Fig.1.

